AD-A139 172 LONG WAVELENGTH LIMIT OF THE CURRENT CONVECTIVE INSTABILITY(U) NAVAL RESEARCH LAB WASHINGTON DC J D HUBA 23 FEB 84 NRL-MR-5264

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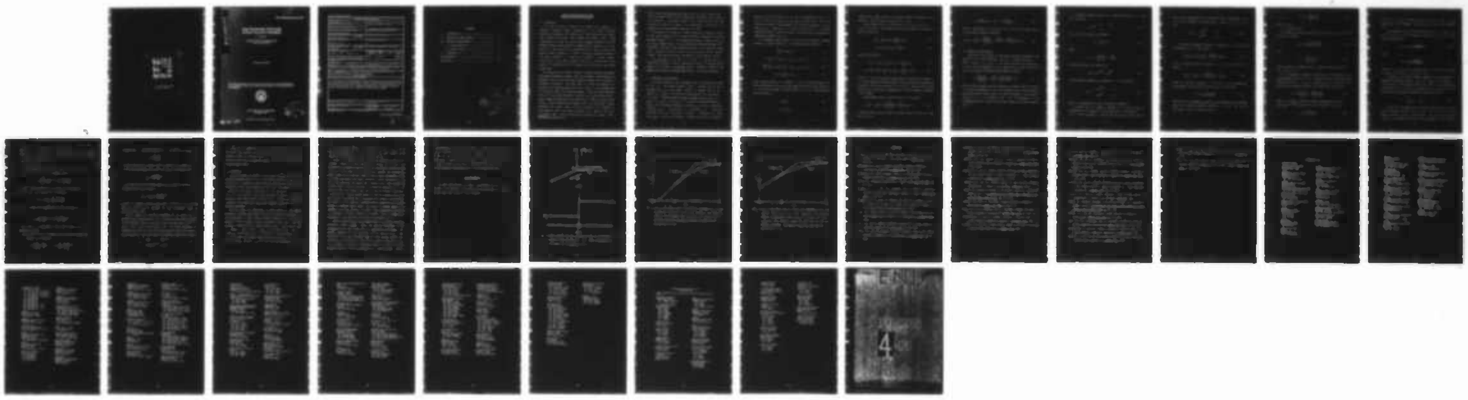

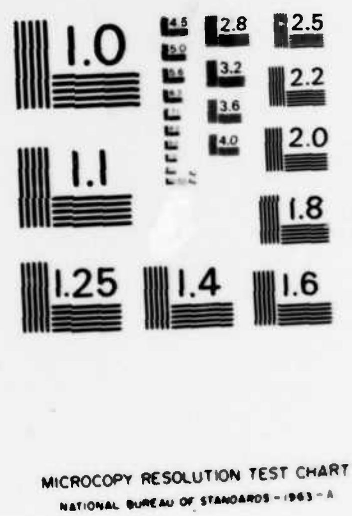
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NRL Memorandum Report 5264

Long Wavelength Limit of the Current Convective Instability

J. D. HUBA

Geophysical and Plasma Dynamics Branch
Plasma Physics Division

February 23, 1984

This research was sponsored by the Defense Nuclear Agency under Subtask S99QMXBC, work unit 00067, work unit title "Plasma Structure Evolution," and by the Office of Naval Research.

NAVAL RESEARCH LABORATORY
Washington, D.C.

REPORT DOCUMENTATION PAGE

1a. REPORT SECURITY CLASSIFICATION UNCLASSIFIED	1b. RESTRICTIVE MARKINGS
2a. SECURITY CLASSIFICATION AUTHORITY	3. DISTRIBUTION/AVAILABILITY OF REPORT Approved for public release; distribution unlimited.
2b. DECLASSIFICATION/DOWNGRADING SCHEDULE	
4. PERFORMING ORGANIZATION REPORT NUMBER(S) NRL Memorandum Report 5264	5. MONITORING ORGANIZATION REPORT NUMBER(S)
6a. NAME OF PERFORMING ORGANIZATION Naval Research Laboratory — 6b. OFFICE SYMBOL (If applicable)	7a. NAME OF MONITORING ORGANIZATION
6c. ADDRESS (City, State and ZIP Code) Washington, DC 20375	7b. ADDRESS (City, State and ZIP Code)
8a. NAME OF FUNDING/SPONSORING ORGANIZATION DNA and ONR — 8b. OFFICE SYMBOL (If applicable)	9. PROCUREMENT INSTRUMENT IDENTIFICATION NUMBER
8c. ADDRESS (City, State and ZIP Code) Washington, DC 20305 Arlington, VA 22217	10. SOURCE OF FUNDING NOS.

PROGRAM ELEMENT NO.	PROJECT NO.	TASK NO.	WORK UNIT NO.
62715H	RR033-02-44		47-0889-0-4 47-0883-0-4

11. TITLE (Include Security Classification) LONG WAVELENGTH LIMIT OF THE CURRENT CONVECTIVE INSTABILITY

12. PERSONAL AUTHOR(S)
J. D. Huba

13a. TYPE OF REPORT	13b. TIME COVERED	14. DATE OF REPORT (Yr., Mo., Day)	15. PAGE COUNT
Interim	FROM ______ TO ______	February 23, 1984	34

16. SUPPLEMENTARY NOTATION This research was sponsored by the Defense Nuclear Agency under Subtask S99QMXBC, work unit 00067, work unit title "*Plasma Structure Evolution*," and by the *Office of Naval Research.*

17. COSATI CODES			18. SUBJECT TERMS (Continue on reverse if necessary and identify by block number)
FIELD	GROUP	SUB. GR.	Current convective instability
			Long wavelength interchange mode
			Auroral ionosphere

19. ABSTRACT (Continue on reverse if necessary and identify by block number)

A linear theory of the current convective instability in the long wavelength limit, i.e., $kL < 1$ where k is the wavenumber and L is the scale length of the density inhomogeneity, is presented. A relatively simple dispersion equation is derived which describes the modes in this limit. Analytical solutions are presented in both the collisional ($\nu_{in} > \omega$) and inertial ($\nu_{in} < \omega$) limits where ω is the wave frequency and ν_{in} is the ion-neutral collision frequency. It is shown that the growth rate scales as k in the collisional limit and as $k^{2/3}$ in the inertial limit. The analytical solutions are compared to exact numerical solutions and very good agreement is found. Applications to the auroral ionosphere are discussed.

20. DISTRIBUTION/AVAILABILITY OF ABSTRACT UNCLASSIFIED/UNLIMITED ☒ SAME AS RPT. ☐ DTIC USERS ☐	21. ABSTRACT SECURITY CLASSIFICATION UNCLASSIFIED
22a. NAME OF RESPONSIBLE INDIVIDUAL J. D. Huba	22b. TELEPHONE NUMBER (Include Area Code) (202) 767-3630 — 22c. OFFICE SYMBOL Code 4780

DD FORM 1473, 83 APR EDITION OF 1 JAN 73 IS OBSOLETE.

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LONG WAVELENGTH LIMIT OF THE CURRENT CONVECTIVE INSTABILITY

I. INTRODUCTION

The current convective instability has recently been suggested as a mechanism to generate density irregularities in the auroral ionosphere (Ossakow and Chaturvedi, 1979; Vickrey et al., 1980; Chaturvedi and Ossakow, 1981; Keskinen and Ossakow, 1982; Keskinen and Ossakow, 1983). These irregularities can cause the scintillation phenomena observed by the DNA Wideband satellite during periods of diffuse aurora (Fremouw et al., 1977; Rino et al., 1978), and will be an important area of study of the DNA HILAT satellite mission (Fremouw et al., 1983). The current convective instability can become unstable in a plasma which supports a density gradient (perpendicular to the ambient magnetic field) and a current which flows parallel to the ambient $\underset{\sim}{B}$ field; a situation which can occur in the diffuse auroral zone.

The instability was initially studied to understand plasma disturbances in laboratory experiments (Lehnert, 1958; Hoh and Lehnert, 1960; Kadomtsev and Nedospasov, 1960) but has been more vigorously pursued lately in regard to ionospheric disturbances. The linear theory of the mode is reasonably well developed in the short wavelength limit ($kL >> 1$ where k is the wavenumber and L is the scale length of the density gradient). Among the issues considered thus far are the linear properties of the mode in the low altitude auroral F region (Ossakow and Chaturvedi, 1979; Vickrey et al., 1980), the high altitude auroral F region (Chaturvedi and Ossakow, 1981), and the auroral E region (Chaturvedi and Ossakow, 1981); as well as studies of the influence of electromagnetic effects (Chaturvedi and Ossakow, 1981), magnetic shear (Huba and Ossakow, 1980), a warm electron beam (Chaturvedi and Ossakow, 1983),

Manuscript approved November 9, 1983.

velocity shear (Satyanarayana and Ossakow, 1983), and kinetic effects (e.g., finite ion Larmor radius effects, wave-particle resonances) (Gary, 1983) on the instability. A nonlinear theory of this instability has also been developed (Chaturvedi and Ossakow, 1979) and numerical simulations of the instability have been performed (Keskinen et al., 1980).

The purpose of this paper is to present analytical and numerical results of the linear properties of the current convective instability in the long wavelength regime ($kL < 1$). The mathematical analysis is similar to that of the long wavelength limit of the $\underset{\sim}{E} \times \underset{\sim}{B}$ instability (from a previous work). ~~(Huba and Zalesak, 1983)~~. The author ~~We~~ derives a relatively simple dispersion equation of the mode, and presents simple dispersion relations in both the ion collisional and ion inertial regimes. These analytical results are compared to exact numerical results.

The organization of the paper is as follows: 1) the derivation of a ~~In the next section we derive the~~ differential equation describing the mode, 2) the presentation of ~~In Section III we present~~ both analytical and numerical results, and 3) a summary of ~~Finally, in Section IV we summarize our~~ findings and a discussion of ~~discuss~~ applications to the auroral ionosphere.

II. DERIVATION OF MODE EQUATION

The plasma configuration and slab geometry used in the analysis are shown in Fig. 1a. The ambient magnetic field is constant and in the z direction $\left(\underset{\sim}{B} = B_0 \hat{e}_z\right)$, the ambient current is constant and in the z direction $\left(\underset{\sim}{J} = J_0 \hat{e}_z\right)$, and the density is inhomogeneous in the x direction $\left(n = n_0(x)\right)$. A weakly collisional plasma is assumed such that $\nu_{ei}/\Omega_e \ll 1$, $\nu_{en}/\Omega_e \ll 1$, $\nu_{ie}/\Omega_i \ll 1$, and $\nu_{in}/\Omega_i \ll 1$ (F region approximation) where $\Omega_\alpha = eB_0/m_\alpha c$ is the cyclotron frequency of species α, ν_{ei} refers to electron-ion collisions, ν_{en} to electron-neutral collisions, ν_{ie} to ion-electron collisions, and ν_{in} to ion-neutral collisions. Furthermore, we assume that $\nu_{en}/\Omega_e \ll \nu_{in}/\Omega_i$ in our analysis. Perturbed quantities are assumed to

vary as $\delta p = \delta p(x) \exp\left[i\left(k_y y + k_z z - \omega t\right)\right]$ and it is assumed that $\omega/\Omega_i \ll 1$, $k\rho_i \ll 1$, and $k_z \ll k_y$, where ρ_i is the mean ion Larmor radius. That is, we consider low frequency, long wavelength, field-aligned perturbations. We also assume $k_z\lambda_{MFP} \ll 1$ where λ_{MFP} is the electron mean free path. We neglect temperature effects. Finally, we consider only electrostatic oscillations and assume quasi-neutrality $\left(n_e = n_i\right)$.

The fundamental equations used in the analysis are continuity, momentum transfer, and charge conservation, in the neutral frame of reference:

$$\frac{\partial n_\alpha}{\partial t} + \nabla\cdot\left(n\underset{\sim}{V}_\alpha\right) = 0 \tag{1}$$

$$0 = -\frac{e}{m_e}\left(\underset{\sim}{E} + \frac{1}{c}\underset{\sim}{V}_e \times \underset{\sim}{B}\right) - \nu_{en}\underset{\sim}{V}_e - \nu_{ei}\left(\underset{\sim}{V}_e - \underset{\sim}{V}_i\right) \tag{2}$$

$$\frac{d\underset{\sim}{V}_i}{dt} = \frac{e}{m_i}\left(\underset{\sim}{E} + \frac{1}{c}\underset{\sim}{V}_i \times \underset{\sim}{B}\right) - \nu_{in}\underset{\sim}{V}_i - \nu_{ie}\left(\underset{\sim}{V}_i - \underset{\sim}{V}_e\right) \tag{3}$$

$$\nabla\cdot\underset{\sim}{J} = \nabla\cdot\left[n_e\left(\underset{\sim}{V}_i - \underset{\sim}{V}_e\right)\right] = 0 \tag{4}$$

where α denotes species (e: electrons, i: ions) and other variables have their usual meaning. Note that we have neglected electron inertia effects in Eq. (2) but have retained ion inertia effects in Eq. (3). The equilibrium drifts are given by

$$\underset{\sim}{V}_{e0} = 0 \tag{5}$$

$$\underset{\sim}{V}_{i0} = V_d\,\hat{e}_z \tag{6}$$

where we have chosen to work in the electron frame of reference in the z direction. Thus, the current is given by $\underset{\sim}{J} = en\, V_d\, \hat{e}_z$.

We now linearize Eqs. (1)-(3) and take $n = n_0 + \delta n$, $\underset{\sim}{V}_\alpha = \underset{\sim}{V}_{\alpha 0} + \delta \underset{\sim}{V}_\alpha$, and $\underset{\sim}{E} = -\nabla\phi$ where ϕ is the perturbed electrostatic potential. Using Eqs. (2) and (3) we find that

$$\delta \underset{\sim}{V}_e = -\frac{c}{B}\nabla_\perp\phi \times \hat{e}_z + \frac{\Omega_e}{\nu_{ei}}\frac{c}{B}\nabla_\parallel\phi\, \hat{e}_z \tag{7}$$

$$\delta \underset{\sim}{V}_i = -\frac{c}{B}\nabla_\perp\phi \times \hat{e}_z + i\,\frac{\tilde{\omega}}{\Omega_i}\nabla_\perp\phi \tag{8}$$

which can be written as

$$\delta \underset{\sim}{V}_e = -i\,\frac{c}{B}k_y\phi\, \hat{e}_x + \frac{c}{B}\phi'\, \hat{e}_y + i\,\frac{c}{B}\frac{\Omega_e}{\nu_{ei}}k_z\phi\, \hat{e}_z \tag{9}$$

$$\delta \underset{\sim}{V}_i = \frac{c}{B}\left[-ik_y\phi + i\,\frac{\tilde{\omega}}{\Omega_i}\phi'\right]\hat{e}_x + \frac{c}{B}\left[-k_y\frac{\tilde{\omega}}{\Omega_i}\phi + \phi'\right]\hat{e}_y. \tag{10}$$

where $\tilde{\omega} = \omega - k_z V_d + i\nu_{in}$ and the prime indicates a derivative with respect to x. We have neglected collisional effects on the electron motion perpendicular to $\underset{\sim}{B}$ and on the ion motion parallel to $\underset{\sim}{B}$. This is justified since we have taken $\nu_{en}/\Omega_e \ll \tilde{\omega}/\Omega_i$ which is appropriate for auroral ionospheric conditions in the F region.

We now substitute Eqs. (9) and (10) into Eq. (4) and obtain

$$(n_0\phi')' - n_0k_y^2\left(1 + i\,\frac{k_z^2}{k_y^2}\frac{\Omega_i\Omega_e}{\tilde{\omega}\,\nu_{ei}}\right)\phi + \frac{B}{c}\frac{\Omega_i}{\tilde{\omega}}k_z V_d\delta n = 0. \tag{11}$$

We relate δn and ϕ using the electron continuity equation and find that (from Eqs. (1) and (9)

$$\delta n = \frac{c}{B}\frac{k_y}{\omega}\left[n_0\phi' - (n_0\phi)' + in_0\frac{k_z}{k_y}\frac{\Omega_e}{\nu_{ei}}k_z\phi\right] \qquad (12)$$

Finally, substutiting Eq. (12) into Eq. (11) we arrive at the mode equation for the current convective instability

$$(n_0\phi') - n_0k_y^2\left[1 + i\frac{k_z^2}{k_y^2}\frac{\Omega_i\Omega_e}{\tilde{\omega}\,\nu_{ei}}\left(1 - \frac{k_zV_d}{\omega}\right)\right]\phi - \frac{k_zV_d}{\omega}\frac{\Omega_i}{\tilde{\omega}}k_yn_0'\phi = 0 \qquad (13)$$

III. ANALYSIS OF MODE EQUATION

The bulk of linear analyses of the current convective instability have made use of the local approximation. That is, it is assumed that $k_y^2L^2 \gg k_x^2L^2 \gg 1$ where $L = (\partial \ln n_0/\partial x)^{-1}$ is the density gradient scale length and $k_x = \partial/\partial x$. With this assumption, one can neglect the first term on the LHS of Eq. (13) and obtain a relatively simple dispersion equation

$$1 + i\frac{k_z^2}{k_y^2}\frac{\Omega_i\Omega_e}{\tilde{\omega}\,\nu_{ei}}\left(1 - \frac{k_zV_d}{\omega}\right) - \frac{k_z}{k_y}\frac{\Omega_i}{\tilde{\omega}}\frac{V_d}{\omega}\frac{n_0'}{n_0}\phi = 0 \qquad (14)$$

which has been thoroughly analyzed (Chaturvedi and Ossakow, 1981). The heart of the local approximation is that the wavelengths of the perturbations are much smaller than the scale length of the density gradient. In this paper we solve Eq. (13) in the opposite limit, i.e., the wavelengths of the perturbations are much larger than the scale length of the density gradient ($kL \ll 1$).

We consider a density profile with a single discontinuity at $x = 0$ (see Fig. 1b) given by

$$n_0(x) = \begin{cases} n_1 & x > 0 \\ n_2 & x < 0 \end{cases} \tag{15}$$

For $x \neq 0$, $n_0' = 0$ and Eq. (13) reduces to

$$\phi'' - k_y^2 \Gamma^2 \phi = 0 \tag{16}$$

where

$$\Gamma^2 = 1 + i \frac{k_z^2}{k_y^2} \frac{\Omega_i \Omega_e}{\tilde{\omega}\, \nu_{ei}} \left(1 - \frac{k_z V_d}{\omega}\right) \tag{17}$$

the solution to Eq. (16) is taken to be

$$\phi(x) = \phi_1 e^{-k_y \Gamma x} + \phi_2 e^{k_y \Gamma x} \tag{18}$$

Since the modes are assumed to be bounded as $x \to \pm\infty$ we note that

$$\phi(x) = \begin{cases} \phi_1 e^{-k_y \Gamma x} & x > 0 \\ \phi_2 e^{k_y \Gamma x} & x < 0 \end{cases} \tag{19}$$

where it is assumed that $\omega >> k_z V_d$ (to be shown a posteriori).

We require that the tangential component of the electric field be continuous at $x = 0$ (Hasegawa, 1971) which means that ϕ is continuous at $x = 0$. This is equivalent to requiring that the interface velocity and the

fluid velocity perpendicular to the interface be equal (Chandresekhar, 1961), i.e., δV_x is continuous at the discontinuity. Thus, $\phi_1 = \phi_2$ in Eq. (19) so that

$$\phi(x) = \begin{cases} \phi_0 e^{-k_y \Gamma x} & x > 0 \\ \phi_0 e^{k_y \Gamma x} & x < 0 \end{cases} \tag{20}$$

To obtain the dispersion equation we integrate Eq. (13) across the discontinuity at x = 0. Thus, we have

$$\int_{-\varepsilon}^{\varepsilon} (n_0 \phi')' dx = \int_{-\varepsilon}^{\varepsilon} \left[n_0 k_y^2 \Gamma^2 \phi - \frac{k_z V_d}{\omega} \frac{\Omega_i}{\tilde{\omega}} k_y n_0' \phi \right] dx \tag{21}$$

Since ϕ is continuous across the boundary at $x = 0$, it is found that Eq. (21) leads to

$$\left(n_0 \phi'\right)_1 - \left(n_0 \phi'\right)_2 = - \frac{k_z V_d}{\omega} \frac{\Omega_i}{\tilde{\omega}} k_y \left(n_1 - n_2\right) \phi_0 \tag{22}$$

where (1,2) indicate the region $x > 0$ ($+\varepsilon$) and $x < 0$ ($-\varepsilon$), respectively. Substituting Eq. (20) into Eq. (22) and letting $\varepsilon \to 0$ we find that

$$\omega \tilde{\omega} \Gamma = k_z V_d \Omega_i \frac{n_1 - n_2}{n_1 + n_2} . \tag{23}$$

Equation (23) is the dispersion equation which describes the long wavelength modes of the current convective instability. From Eq. (23) we note that $\omega / k_z V_d \sim \Omega_i / \tilde{\omega} >> 1$ so that we can take

$$\Gamma^2 = 1 + i\,\frac{k_z^2}{k_y^2}\,\frac{\Omega_i}{\tilde{\omega}}\,\frac{\Omega_e}{\nu_{ei}} \tag{24}$$

and $\tilde{\omega} = \omega + i\nu_{in}$.

A. Collisional Limit

We consider the collisional limit $\nu_{in} >> \omega$ so that $\tilde{\omega} = i\nu_{in}$ in Eqs. (23) and (24). We find that

$$\gamma_{n\ell} = -\,k_z V_d\,\frac{\Omega_i}{\nu_{in}}\,\frac{1}{\Gamma}\,\frac{n_1 - n_2}{n_1 + n_2} \tag{25}$$

with

$$\Gamma = \left(1 + \frac{k_z^2}{k_y^2}\,\frac{\Omega_i}{\nu_{in}}\,\frac{\Omega_e}{\nu_{ei}}\right)^{1/2} \tag{26}$$

where the subscript $n\ell$ denotes nonlocal. Instability can occur for $k_z V_d$ $(n_1 - n_2)/(n_1 + n_2) < 0$.

It is interesting to compare the growth rate of the instability in the long wavelength limit ($kL << 1$, i.e., nonlocal) to that of the short wavelength limit ($kL >> 1$, i.e., local). Assuming $\nu_{in} >> \omega$, we find from Eq. (14) that the short wavelength growth rate (local growth rate γ_ℓ) is given by

$$\gamma_\ell = -\,\frac{k_z}{k_y}\,\frac{\Omega_i}{\nu_{in}}\,\frac{V_d}{L}\left(1 + \frac{k_z^2}{k_y^2}\,\frac{\Omega_i}{\nu_{in}}\,\frac{\Omega_e}{\nu_{ei}}\right)^{-1} \tag{27}$$

where $L = \left(n_0'/n_0\right)^{-1}_{max}$. Defining an effective wavenumber $\hat{k} = \left(k_y^2 + k_z^2\right.$ $\left.\left(\Omega_e\Omega_i/\nu_{in}\nu_{ei}\right)\right)^{1/2}$ we can rewrite Eq. (25) in terms of Eq. (27), i.e.,

$$\gamma_{n\ell} = \hat{k}L\,\frac{n_1 - n_2}{n_1 + n_2}\,\gamma_\ell \tag{28}$$

For $n_1 >> n_2$ we have the simple relation $\gamma_{n\ell} = \hat{k}L\,\gamma_\ell$ so that (1) $\gamma_{n\ell}$ is proportional to the "wavenumber" $\hat{k}$ and (2) $\gamma_{n\ell} << \gamma_\ell$ since we have assumed $\hat{k}L << 1$.

We now compare the analytical expressions derived for the growth rates (Eqs. (25) and (27)) with the numerical solution of Eq. (13). We choose a density profile given by

$$n(x) = n_0 \frac{1 + \varepsilon \tanh(x/a)}{1 - \varepsilon} \tag{29}$$

so that

$$\frac{n'}{n} = \frac{1}{a} \frac{\varepsilon \operatorname{sech}^2(x/a)}{1 + \varepsilon \tanh(x/a)} \tag{30}$$

We take $\varepsilon = 0.8$ so that n'/n is a maximum at $x/a = -0.55$. We find then that $(n'/n)_{max} = (1/a)$ so that $L \simeq a$ in Eqs. (27)-(30).

We consider the following physical parameters for the comparison of analytical and numerical results: $\nu_{ei}/\Omega_e = 10^{-4}$, $\nu_{in}/\Omega_i = 10^{-4}$ and $k_z/k_y = 10^{-4}$ (Rino et al., 1978; Ossakow and Chaturvedi, 1979). We also consider the normalization $\hat{\omega} \equiv \omega|L/V_d|$ where it is assumed that $L/V_d < 0$ so that $\gamma > 0$. We find then that the analytical expression for the growth rate is

$$\hat{\gamma} = \begin{cases} 0.50 & kL >> 1 \\ 0.56\, k_y L & kL << 1 \end{cases} \tag{31}$$

We solve Eq. (13) numerically for the density profile given by Eq. (29) and the parameters described above. In Fig. (2) we plot $\hat{\gamma}$ vs. $k_y L$. The numerical solution is plotted in the regime $0.1 < k_y L < 10.0$ and is labelled 'exact'. As is evident, the numerical solution asymptotes to the analytical

expression (Eq. (31)) in both the short wavelength ($kL >> 1$) and long wavelength ($k << 1$) limits. Thus, Eqs. (25) and (27) provide good estimates of the current convective instability in the long wavelength and short wavelength limits, respectively.

B. Inertial Limit

We now consider the ion inertial limit given by $\omega >> \nu_{in}$ so that $\tilde{\omega} = \omega$ in Eqs. (23) and (24). The dispersion equation is

$$\omega^2 \left(1 + i \frac{k_z^2}{k_y^2} \frac{\Omega_i}{\omega} \frac{\Omega_e}{\nu_{ei}}\right)^{1/2} = k_z V_d \Omega_i \frac{n_1 - n_2}{n_1 + n_2} \tag{32}$$

We obtain a simple expression for the growth rate by assuming that $\omega/\Omega_i <<$ $(k_y^2/k_y^2)(\Omega_e/\nu_{ei})$. In this limit Eq. (32) becomes

$$\omega^3 = - i k_y^2 V_d^2 \Omega_i \frac{\nu_{ei}}{\Omega_e} \left(\frac{n_1 - n_2}{n_1 + n_2}\right)^2 \tag{33}$$

Again, we consider the normalization $\hat{\omega} = \omega |L/V_d|$ and take $L/V_d < 0$. We find that

$$\hat{\gamma}_{n\ell} = \left(\frac{\nu_{ei}}{\Omega_e} \frac{\Omega_i L}{V_d}\right)^{1/3} (k_y L)^{2/3} \left(\frac{n_1 - n_2}{n_1 + n_2}\right)^{2/3} \tag{34}$$

In this limit it is interesting to note that $\hat{\gamma}$ scales as $(k_y L)^{2/3}$ and is independent of k_z/k_y.

We contrast Eq. (34) to the growth rate in the short wavelength limit ($kL >> 1$). From Eq. (14) we find that

$$\hat{\gamma}_\ell = - \frac{1}{2} \frac{k_z^2}{k_y^2} \frac{\Omega_e}{\nu_{ei}} \left|\frac{L\Omega_i}{V_d}\right| \left[1 - \left(1 + 4 \left|\frac{V_d}{L\Omega_i}\right| \frac{k_y^3}{k_z^3} \frac{\nu_{ei}^2}{\Omega_e^2}\right)^{1/2}\right] \tag{35}$$

Assuming that $4|V_d/L\Omega_i|(k_y^3/k_z^3)(\nu_{ei}^2/\Omega_e^2) \ll 1$ we find that Eq. (36) reduces to

$$\hat{\gamma}_\ell = \frac{k_y}{k_z}\frac{\nu_{ei}}{\Omega_e} \tag{36}$$

Following Chaturvedi and Ossakow (1981) we maximize $\hat{\gamma}_\ell$ with respect to k_z/k_y and find that the maximum local growth rate $(\hat{\gamma}_{\ell m})$ in the inertial limit is

$$\hat{\gamma}_{\ell m} = \left(\frac{\nu_{ei}}{\Omega_e}\frac{L\Omega_i}{4V_d}\right)^{1/3} \tag{37}$$

We can express the long wavelength growth rate in terms of the maximum short wavelength growth rate to obtain

$$\hat{\gamma}_{n\ell} \simeq \alpha\gamma_{\ell m}\left[k_y L \frac{n_1 - n_2}{n_1 + n_2}\right]^{2/3} \tag{38}$$

where α is a numerical factor of order unity. Thus, for $n_1 \gg n_2$ we find that $\gamma_{n\ell} \simeq (k_y L)^{2/3}\, \gamma_{\ell m}$, in contrast to the collisional limit in which $\hat{\gamma}_{n\ell}$ scales as $k_y L$.

We now compare the analytical expressions for the growth rate with numerical results based upon Eq. (13). We choose the same density profile as in the collisional case (Eq. (29) with $\varepsilon = 0.8$), and take $\nu_{ei}/\Omega_e = 10^{-4}$, $k_z/k_y = 10^{-4}$, $\nu_{in}/\Omega_i = 0$, and $|L\Omega_i/V_d| = 10^4$. For these parameters, the assumption that led to Eq. (36) is not satisfied and the analytical expression for $\hat{\gamma}_\ell$ given by Eq. (35) must be used. The analytical growth rate is given by

$$\hat{\gamma} = \begin{cases} 0.62 & k_y L \gg 1 \\ 0.86(k_y L)^{2/3} & k_y L \ll 1 \end{cases} \tag{39}$$

The results of the comparison are shown in Fig. 3 where we plot $\hat{\gamma}$ vs. k_yL. The growth rate $\hat{\gamma}$ given by Eq. (39) is plotted as shown on Fig. (3) and the numerical results are labelled 'exact.' Again, excellent agreement is found between the analytical and numerical results in both the short wavelength and long wavelength limits.

IV. DISCUSSION

We have presented an analysis of the current convective instability in the long wavelength limit ($kL \ll 1$). The principal result of the paper is the derivation of a relatively simple dispersion equation which describes the instability in the long wavelength limit (Eq. (22)). This equation is solved analytically in two limits: collisional $(\nu_{in} \gg \omega)$ and inertial $(\nu_{in} \ll \omega)$. We find that in the collisional limit the growth rate scales as k, while in the inertial limit it scales as $k^{2/3}$. We have also presented a comparison of the analytical results with numerical results and have found very good agreement.

We now discuss the application of these results to the auroral ionosphere. Chaturvedi and Ossakow (1981) have discussed the relevance of the short wavelength (or local) current convective instability to both the low altitude (~ 400 km) auroral F region (collisional limit) and to the high altitude (~ 1000 km) auroral F region (inertial limit). For the low altitude F region they find that $\gamma_\ell \sim 3 \times 10^{-3}$ sec^{-1}, for $V_d \sim 500$ m/sec, $\nu_{ei} \sim 5 \times 10^2$ sec^{-1}, $\nu_{in} \sim 5 \times 10^{-2}$ sec^{-1}, $L \sim 50$ km and $m_e/m_i \sim 3 \times 10^{-5}$; for the high altitude F region they also find that $\gamma_\ell \sim 3 \times 10^{-3}$ sec^{-1} but for $\nu_{ei} \sim 30$ sec^{-1}, $\nu_{in} \lesssim 10^{-3}$ sec^{-1} and other parameters the same

as the low altitude F region. From our analysis we would predict that the long wavelength modes ($kL < 1$) would have much longer growth times ($1/\gamma$), particularly in the low altitude F region.

There is a difficulty in applying the long wavelength results to the auroral F region for large density scale lengths ($L \sim 50$ km). The perpendicular wavelengths associated with the long wavelength modes are such that $k_\perp L < 1$ which leads to $\lambda_\perp > 2\pi L \sim 300$ km. However, the perpendicular spatial scale size of observed auroral F region density enhancements, which can provide the zeroth order density gradients to drive the instability, can be of this magnitude (few hundred kms) (Vickrey et al., 1980; Tsunoda and Vickrey, 1983) so that it is difficult to satisfy the condition $k_\perp L < 1$. On the other hand, much smaller scale sizes of auroral structure occur during discrete aurora. Discrete aurora appear to have two distinct scale sizes. One is tens of kilometers and is associated with inverted V precipitation. The other is of the order of a kilometer and is associated with discrete auroral arc elements (Davis, 1978). The auroral arc elements appear to be imbedded in the larger inverted V structure. Thus, application of our theory to structure in discrete auroral arc elements leads to perpendicular wavelengths $\lambda_\perp \gtrsim$ few kms which can satisfy the requirement $k_\perp L < 1$.

Finally, several aspects of the present theory of the current convective instability deserve comment. First, the parallel wavelengths of the irregularities are much larger than the perpendicular wavelengths. For typical parameters we find that $k_\parallel / k_\perp \sim 10^{-4}$ so that $\lambda_\parallel \sim 10^4 \lambda_\perp$. The instability in the long wavelength limit can then produce parallel wave structures $\lambda_\parallel \gtrsim 10^4$ km which is larger than the parallel system size. Thus, it is important to develop an appropriate theory to consider the finite parallel extent of the auroral ionosphere (e.g., Sperling (1983)). Second,

electromagnetic effects also need to be considered in a more detailed theory. Chaturvedi and Ossakow (1981) investigated these effects on the local theory of the current convective instability and found them to be negligible. However, the electromagnetic corrections are proportional to ω_{pe}^2/c^2k^2 and may become large in the long wavelength limit since $kL < 1$. We are presently investigating both of the above mentioned effects.

Acknowledgments

We thank Mike Keskinen for helpful conversations and for a critical reading of the manuscript. This research has been supported by the Defense Nuclear Agency and the Office of Naval Research.

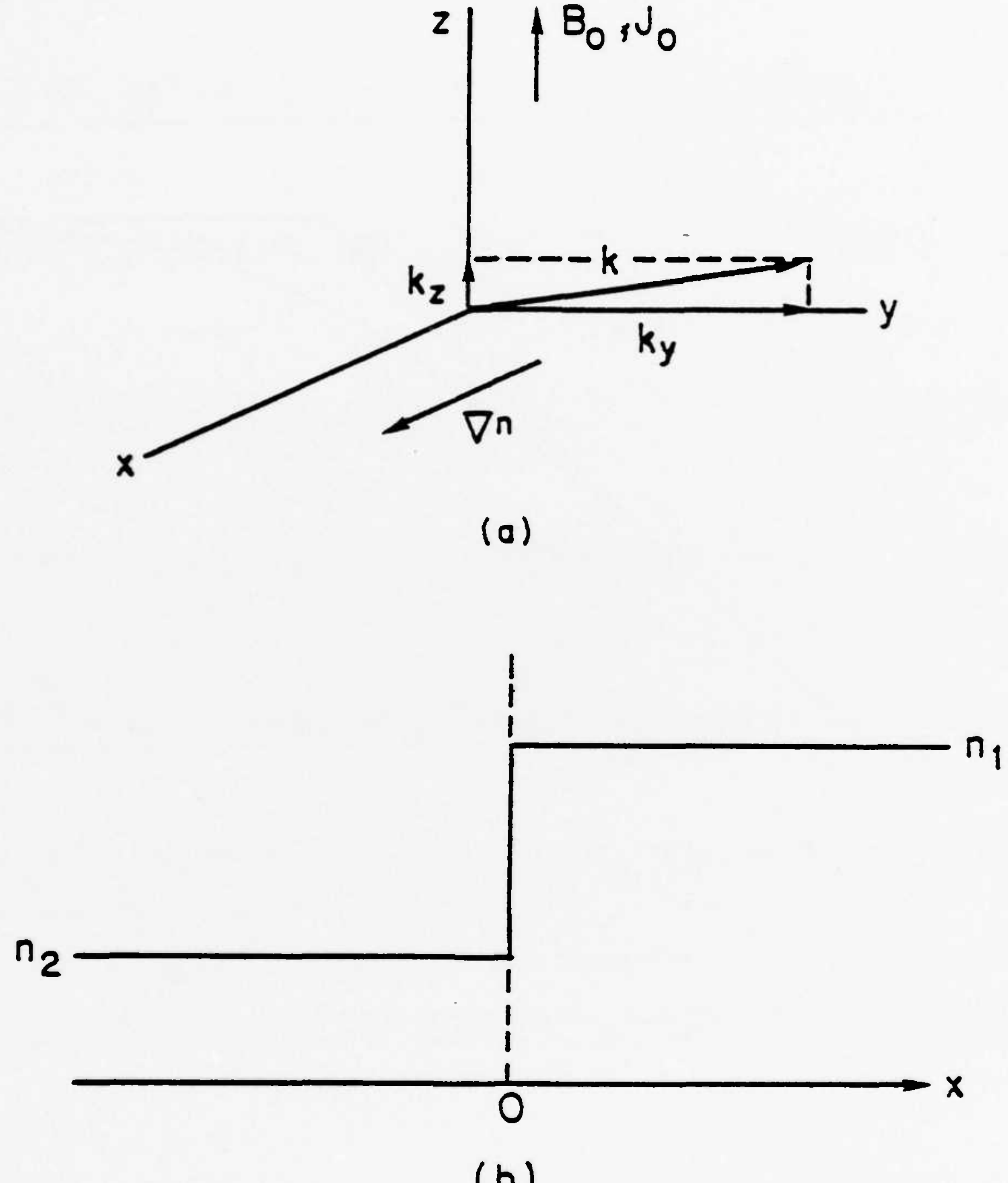

Fig. 1 Plasma geometry and slab configuration used in the analysis. (a) Standard plasma configuration. (b) Plasma configuration with a discontinuity in the density of $x = 0$.

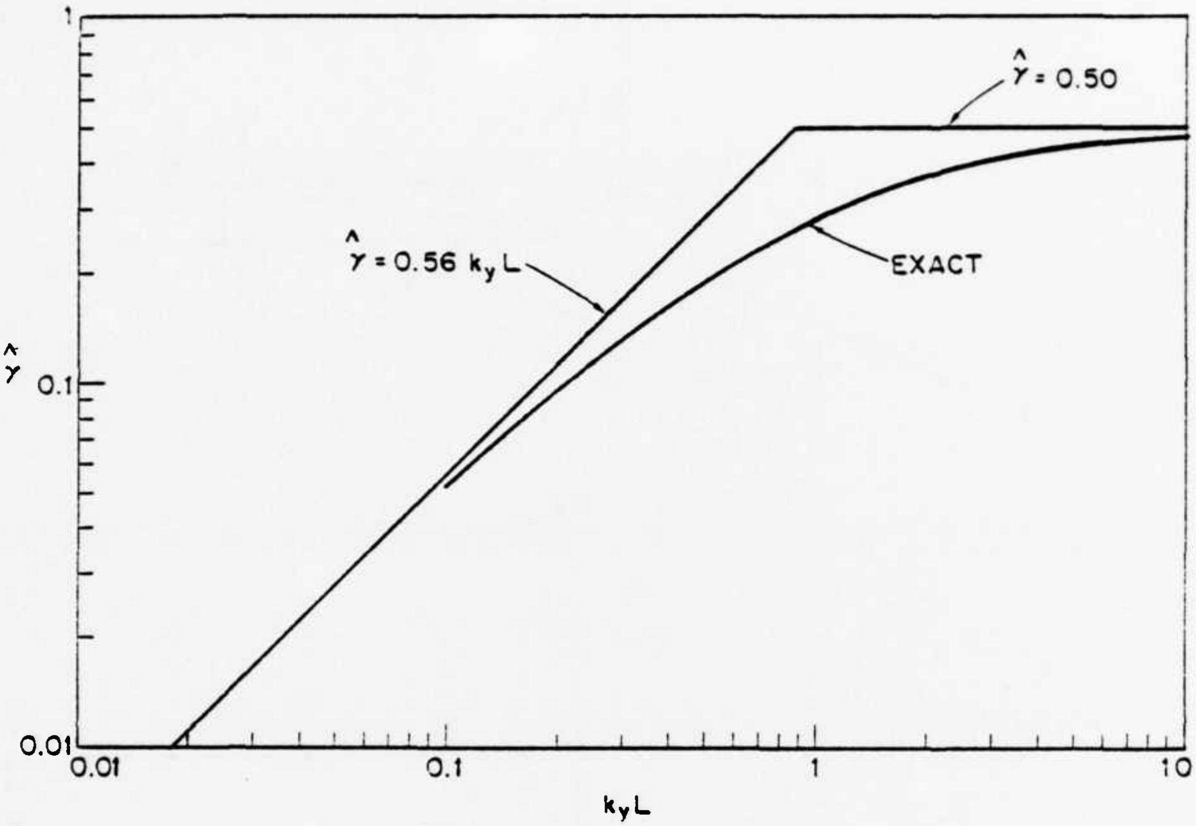

Fig. 2 Plot of the growth rate $(\tilde{\gamma} = \gamma |L/V_d|)$ vs. wavenumber $(k_y L)$ in the collisional limit $(\nu_{in} >> \omega)$ for both analytical and numerical results. The parameters used are $\nu_{ei}/\Omega_e = 10^{-4}$, $\nu_{in}/\Omega_i = 10^{-4}$, and labelled accordingly. The numerical results are based upon Eqs. (13) and (29) with $\varepsilon = 0.8$.

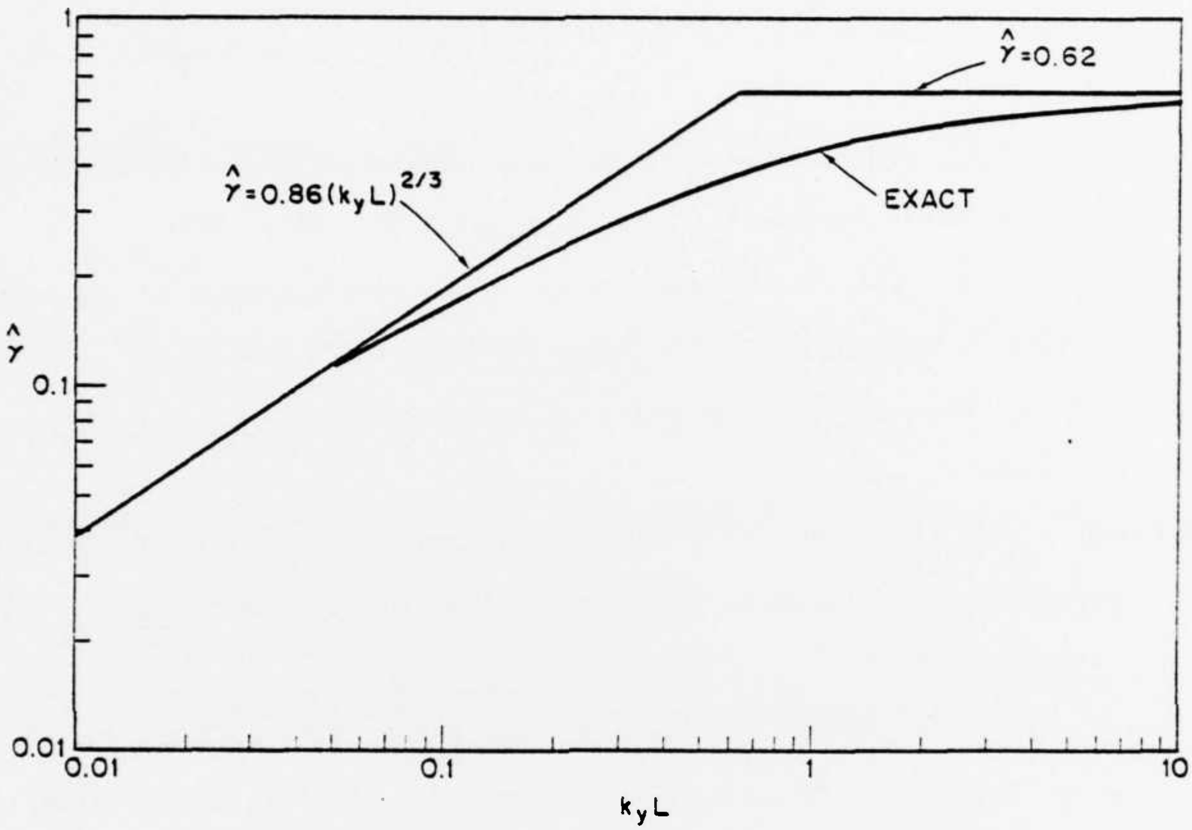

Fig. 3 Plot of the growth rate $\left(\tilde{\gamma} = \gamma|L/V_d|\right)$ vs. wavenumber $\left(k_yL\right)$ in the inertial limit $\left(\nu_{in} << \omega\right)$ for both analytical and numerical results. The parameters used are $\nu_{ei}/\Omega_e = 10^{-4}$, $k_z/k_y = 10^{-4}$, and $\nu_{in}/\Omega_i = 0$. The analytical results are based upon Eq. (39) and labelled accordingly. The numerical results are based upon Eqs. (13) and (29) with $\varepsilon = 0.8$.

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